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Baryon Number Violation via Majorana Neutrinos in the Early Universe, at the LHC, and Deep Underground

Hooman Davoudiasl[1] and Yue Zhang[2]

[1]*Department of Physics, Brookhaven National Laboratory, Upton, NY 11973, USA*
[2]*Walter Burke Institute for Theoretical Physics, California Institute of Technology, Pasadena, CA 91125, USA*

`hooman@bnl.gov` `yuezhang@theory.caltech.edu`

We propose and investigate a novel, minimal, and experimentally testable framework for baryogenesis, dubbed *dexiogenesis*, using baryon number violating effective interactions of right-handed Majorana neutrinos responsible for the seesaw mechanism. The distinct LHC signature of our framework is same-sign top quark final states, possibly originating from displaced vertices. The region of parameters relevant for LHC phenomenology can also yield concomitant signals in nucleon decay experiments. We provide a simple ultraviolet origin for our effective operators, by adding a color-triplet scalar, which could ultimately arise from a grand unified theory.

Gauge singlet right-handed neutrinos (RHNs) provide perhaps the simplest explanation of non-zero neutrino masses, as demanded by a large and well-established body of neutrino oscillation data, thereby allowing the Standard Model (SM) to be a renormalizable theory of Nature. If neutrinos are Dirac particles, the associated RHNs only couple to the SM via Yukawa interactions of negligible strength, $y_N \lesssim 10^{-12}$, and are not expected to be directly detectable in the foreseeable future. On the other hand, if the observed neutrinos are Majorana states, they most naturally get their mass from a seesaw mechanism [1]. In that case, RHNs may have $\mathcal{O}(1)$ couplings to SM neutrinos, but they would then have to be exceedingly heavy, $\gtrsim 10^{14}\,$GeV, far beyond the reach of terrestrial experiments. However, it may very well be that RHNs are much lighter, and they can be directly probed in lepton number violating processes in collider experiments [2–6] and perhaps other searches, such as those for proton decay [7].

In this letter, we assume that the RHNs associated with the seesaw mechanism are near the weak scale $\lesssim 1\,$TeV. Higher dimensional operators involving the RHNs are generically present. In particular, we will further assume that the RHNs have baryon number violating interactions, mediated by dimension 6 operators involving right-handed quarks and suppressed by a scale $\sim 1-10$ TeV. We will show that these assumptions allow for direct generation of a baryon number asymmetry through RHN decays in the early Universe, which we dub *dexiogenesis* (dexios: Greek for the right hand). This is in contrast to canonical leptogenesis [8] where the lepton asymmetry needs to be further processed into baryon number through electroweak sphalerons [9]. Our direct baryogenesis mechanism is most constrained by nucleon decay bounds. However, we show that for viable parameters one could have distinct collider signatures. This scenario, with dim-6 operators, is an effective field theory and can be embedded in a simple renormalizable model, and possibly a grand unified theory (GUT), where additional signals are expected to arise at the Large Hadron Collider (LHC) and future colliders. For a partial list of other works whose subjects have some overlap with that of this letter, see, for example, Refs. [14–17].

Our point of departure is the SM augmented by two Majorana RHNs N_a, $a = 1, 2$, the minimum required for a realistic seesaw mechanism based on current neutrino data. We add the following terms to the Lagrangian

$$\mathcal{L}_N = M_a \bar{N}^c_a N_a + y_N^{ai} H \bar{N}_a L_i + \text{h.c.}\,, \tag{1}$$

where M_a is the mass of N_a, $i = 1, 2, 3$ enumerates SM generations, and y^{ai} is a Yukawa matrix; H and L_i are the Higgs and lepton doublets of the SM, respectively.

Light neutrino masses $m_\nu \lesssim 0.1$ eV, implied by the oscillation data, can be generated from the renormalizable interactions in Eq. (1), via the seesaw mechanism: $(m_\nu)_{ij} \sim y_N^{ai} y_N^{aj} \langle H\rangle^2/M_a$. Nonetheless, the SM, henceforth defined to include the Lagrangian in Eq. (1), is widely expected to be an effective theory that is further enriched with new interactions at higher scales. This expectation is strongly motivated by the need for a dark matter candidate and also a baryogenesis mechanism to generate the observed baryon asymmetry of the Universe.

Baryogenesis requires a source of baryon number violation [18]. Hence, it may be necessary to extend the SM by effective operators that violate baryon number [19]. Such operators are suppressed by scales associated with new physics, often considered to be very high, $\gtrsim 10^{15}$ GeV, as implied by nucleon decay constraints. However, it is compelling to look for scenarios where new physics arises at lower scales. For one thing, it is reasonable to assume that the physics underlying the Higgs potential is not very far from the weak scale. Such physics would then have the added benefit of being potentially testable.

Motivated by the above considerations, we assume the following baryon-number violating operators involving the RHNs, in addition to those made up of only the observed SM fields,

$$\mathcal{L}_{\rm BV} = \frac{\lambda_a^{ijk}}{\Lambda^2}[N_a u_i d_j d_k]_R + \frac{\kappa_a^{ilm}}{\Lambda^2}[N_a d_i]_R [Q_l Q_m]_L + \text{h.c.}\,, \tag{2}$$

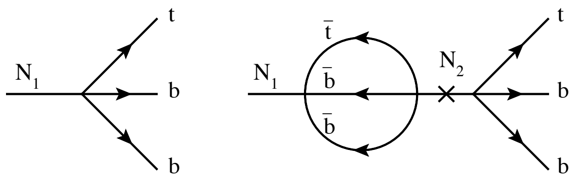

FIG. 1: Tree and two-loop diagrams for dexiogenesis.

where i, j, k are family numbers of right-handed quark (u, d) mass eigenstates and l, m enumerate left-handed quark (Q) generations. Here, λ_a^{ijk} and κ_a^{ilm} are generally complex constants determined by the ultraviolet (UV) theory. These operators could arise from grand unified theories, as shown in a concrete example in the end. They are the lowest dimensional operators that allow RHNs to couple to baryon number directly [20].

In order to have successful dexiogenesis, the coefficients of the relevant dim-6 operators cannot be too small. To avoid excessive low energy baryon number violation for $\Lambda \gtrsim 1$ TeV, those operators would mainly involve right-handed third generation quarks, which help avoid severe constraints from nucleon decay data, discussed in more detail below. To see this, note that quark mass diagonalization can induce operators that involve light quarks, in the presence of left-handed fields. For this reason, we will not further consider tree-level dim-6 operators $NdQQ$ in Eq. (2). These operators can still be generated from $Ntbb$ via radiative corrections, but would not lead to severe constraints.

In light of the above discussion, throughout this work, we will focus on operators involving right-handed third generation top and bottom quarks (originating from the first term in Eq. (2), with explicit spinor contractions)

$$\mathcal{L}_{\rm BV}^{\rm 3R} = \lambda_a \frac{[\bar{N}_a^c P_R b][\bar{t}^c P_R b]}{\Lambda^2}\,, \tag{3}$$

where $P_R \equiv (1+\gamma_5)/2$ is the right-handed projector. The dominance of the third generation could be expected from a connection to UV flavor dynamics. Operators with other combinations of chiralities and flavors can in principle be present, but they must be more suppressed [23].

Remarkably, the addition of the operators in Eq. (3) provides all the necessary ingredients encoded in Sakharov's conditions [18] for baryogenesis: *(i)* these interactions are manifestly baryon number violating, *(ii)* their complex coefficients provide a source of CP violation, and *(iii)* if the Universe has a low reheat temperature $T_{\rm RH} \ll M_a$, then the N_a, assumed to be non-thermally produced throughout this work, will decay out of equilibrium. This mechanism, dexiogenesis, allows $T_{\rm RH} \ll 100$ GeV, since the baryon asymmetry is directly generated and hence electroweak sphalerons do not need to be active.

Let N_1 be the lighter of the two RHNs in our setup. Then, the interference of the tree and the 2-loop diagrams in Fig. 1 will lead to a baryon asymmetry $\varepsilon \equiv \Gamma(N_1 \to tbb) - \Gamma(N_1 \to \bar{t}\bar{b}\bar{b})/(2\Gamma_{N_1})$, where the width of N_1 is given by

$$\Gamma_{N_1} = \frac{|\lambda_1|^2 M_1^5}{1024\pi^3\Lambda^4} F\left(m_t^2/M_1^2\right)\,, \tag{4}$$

with $F(x) = 1 - 8x - 12x^2 \log x + 8x^3 - x^4$.

In the presence of the higher dimensional operator Eq. (3) with a TeV scale cutoff, N_1 decays induced by neutrino Yukawa interactions [Eq. (1)] are subdominant, for values of M_1 near the weak scale. Given a realistic seesaw mechanism for the SM active neutrino masses, in general we have $y_N^a \lesssim 10^{-6}\sqrt{M_1/(200\,\text{GeV})}$ in the absence of fine tuning [24]. The induced $N_1 \to W\ell$ decay rate is then estimated to be $\Gamma_{N_1\to W\ell} \lesssim 10^{-12}\,\text{GeV}\,(y_N^a/10^{-6})^2\,[M_1/(200\,\text{GeV})]$. We find that, for M_1 of a few hundred GeV and $\Lambda/\sqrt{\lambda_1} \lesssim 25\,$TeV, that rate is smaller than the baryonic decay rate.

The baryon asymmetry can be conveniently obtained using the unitarity cut method [25]

$$\varepsilon = \frac{\text{Im}(\lambda_1^2\lambda_2^{*2})}{3072\pi^3|\lambda_1|^2}\left(\frac{M_1}{\Lambda}\right)^4 \frac{M_1 M_2}{(M_2^2 - M_1^2)}\,. \tag{5}$$

The relation between the above asymmetry and the baryon number to entropy ratio $\eta \equiv n_B/s \sim 10^{-10}$ [27] depends on the non-thermal production mechanism for N_1, but it can plausibly be $\eta \sim \varepsilon/100$. For example, let us assume that a heavy modulus, such as an inflaton, decays equally into radiation and N_1, which promptly decays. We will take the reheat temperature to be $T_{\rm RH} \sim 1$ GeV. Then, one can estimate $\eta \sim \varepsilon/g^*$ where $g^* \sim 100$ is the number of relativistic degrees of freedom at $T_{\rm RH}$. Alternatively, if the modulus decays exclusively into N_1, and it is the decay of the N_1 population that reheats the Universe, we end up with $\eta \sim \varepsilon\, T_{\rm RH}/M_1$, which for $M_1 \sim 100$ GeV, again yields $\eta \sim \varepsilon/100$. Hence, for $M_1 \sim M_2$ and $\lambda_a \sim 1$, we typically require $M_1/\Lambda \gtrsim 0.1$. Consequently, for $M_1 \lesssim 1$ TeV, relevant for collider phenomenology, the cutoff scale must be sufficiently low, $\Lambda \lesssim 10$ TeV. Let us then examine the experimental constraints on Λ.

It turns out that nucleon decay limits provide the most stringent lower bound on Λ in the above model where RHNs violate both lepton and baryon numbers. While

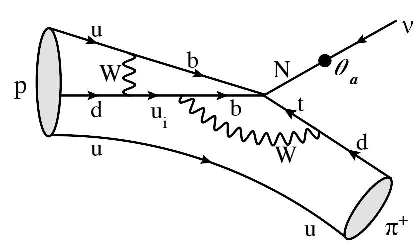

FIG. 2: One of the leading diagrams that yield proton decay.

the operators in Eq. (3) do not contain light quarks, quantum loop corrections can induce nucleon decay via these baryon number violating interactions. Fig. 2 provides a sample two-loop diagram for proton decay, with a rate given by

$$\Gamma(p \to \pi^+ \nu) = \frac{(1+g_A)^2 \alpha^2 m_p}{32\pi f_\pi^2} |\xi|^2 \ , \tag{6}$$

where $g_A = 1.27$ is the nucleon axial charge, $f_\pi = 131\,\mathrm{MeV}$ is pion decay constant, lattice calculations [26] yield the form factor $\alpha \approx -0.01125\,\mathrm{GeV}^3$, and

$$\xi \approx \frac{\Lambda_{qcd} G_F^2 m_t m_b^2 V_{td}^2 V_{ub}^* V_{tb}^*}{(16\pi^2)^2 \Lambda^2} \lambda_a \theta_a \tag{7}$$

is the Wilson coefficient from estimating the two-loop diagram in Fig. 2. The angle θ_a is the mixing between N_a and the SM active neutrinos. The hadronic mass scale $\Lambda_{qcd} \approx 200\,\mathrm{MeV}$ must be introduced under a symmetry argument. The operator we started with is $[\bar{N}_a^c P_R b][\bar{t}^c P_R b]$, and after the W-loop dressing as in Fig. 2, the operator for proton decay turns out to be $[\bar{N}_a^c P_R d][\bar{u}^c P_L d]$ (which is the radiatively generated $NdQQ$ operator mentioned earlier). The fact that one of the down quark is still right-handed implies an external (constituent) quark mass insertion $\sim \Lambda_{qcd}$.

The resulting proton decay life time is

$$\tau(p \to \pi^+ \nu) \approx 2.5 \times 10^{32}\,\mathrm{yr} \left(\frac{\Lambda/\sqrt{\lambda_a}}{1.5\,\mathrm{TeV}}\right)^4 \left(\frac{\theta_a}{10^{-6}}\right)^{-2} \ . \tag{8}$$

The current experimental lower limit on the $p \to \pi^+ \nu$ decay channel is 1.6×10^{31} yr [27]. Hence, the above lifetime (8) is not far from the current limit and, in the region of parameters considered in our work, can be within the reach of future nucleon decay experiments [28, 29].

Here, we also address a potential bound from requiring that the asymmetry in Eq. (5) is not washed out after baryogenesis. The low reheat temperature assumption mentioned before can ensure that processes mediated by the operators in Eq. (3), such as $bb \to N\bar{t}$, are effectively turned off. However, loop processes similar to those depicted in Fig. 2 can lead to baryon number violation mediated by lighter states. Let us assume, for illustrative purposes, that the reheat temperature is $T_{\rm RH} \sim 1$ GeV, well above the temperature at the onset of Big Bang Nucleosynthesis. We then have to make sure that the analog of neutron-anti-neutron oscillation mediated by $(css)(\bar{c}\bar{s}\bar{s})/\Lambda'^5$ is not active. A straightforward comparison with Eq. (7) yields

$$\Lambda'^5 \approx \left(\frac{\theta_a}{\xi} \frac{\Lambda_{qcd}}{m_s} \frac{V_{td}^2 V_{ub}^*}{V_{ts}^2 V_{cb}^*}\right)^2 M_1 \ , \tag{9}$$

where $m_s \approx 100$ MeV is the strange quark mass. For $M_1 = 200$ GeV, we find $\Lambda' \sim 2 \times 10^9$ GeV. The rate of the $css \leftrightarrow \bar{c}\bar{s}\bar{s}$ process is estimated to be $\Gamma_{\Delta B=2} \sim T_{\rm RH}^{11}/\Lambda'^{10} \sim 10^{-93}$ GeV, for $T_{\rm RH} \sim 1$ GeV, which is completely negligible compared to the Hubble rate at this temperature, $H \sim T_{\rm RH}^2/M_{\rm planck} \sim 10^{-19}$ GeV. We also note that bounds from neutron-anti-neutron oscillation would not constrain our model, since that process involves up and down quarks, for which the corresponding suppression scale is larger than the above Λ' scale, and much higher than the current limit [30].

An immediate consequence of Eq. (3) is the possible production of same-sign top quarks at the LHC and future hadron colliders, due to the Majorana nature of RHNs, as shown in the left panel of Fig. 3. In this process, the RHN N_a and a top quark are first produced, and then N_a decays into another top quark and two bottom quarks. Because it is a Majorana particle, an on-shell N_a is equally likely to decay into tbb or $\bar{t}\bar{b}\bar{b}$ final states. The violation of baryon number is manifested in terms of the violation of top quark number (by two units). The sign of the top quark can be inferred from its leptonic decays. For a RHN with a few hundred GeV mass and the effective cutoff scale $\Lambda/\sqrt{\lambda_a}$ of a few TeV, we find that the cross section for this process can be as large as ~ 0.3 fb in the LHC Run-II at 13 TeV. The main background for this signal is from $t\bar{t}b\bar{b}$ final states with the lepton charge from a top quark decay misidentified, which is suppressed by the small misidentification rate [31]. In Table I, we list the leading order cross sections of our signal for several sample mass values of RHNs. These points have not been excluded by the existing LHC data. For example, with $M_a = 200$ GeV and $\Lambda/\sqrt{\lambda_a} = 1.5$ TeV, the cross section at 8 TeV is 0.07 fb, which implies only 1-2 events given the existing integrated luminosity $\sim 27\,\mathrm{fb}^{-1}$, and it is further suppressed by the top quark leptonic branching ratios.

$\sigma(pp \to tN \to ttbb)$				
M_a	200 GeV	500 GeV	800 GeV	1 TeV
$\sqrt{s} = 13$ TeV	0.34 fb	0.16 fb	8×10^{-2} fb	5×10^{-2} fb

TABLE I: Same-sign top quark production cross section, at the 13 TeV LHC, via a Majorana RHN and the contact operators in Eq. (3). The cutoff scale is fixed to be $\Lambda/\sqrt{\lambda_a} = 1.5$ TeV.

Following the same logic as introducing RHNs to make the SM renormalizable, we now discuss a UV completion that generates the effective operator Eq. (3). Given a TeV scale cutoff, it is possible to directly probe the heavy particles in such a model in LHC Run-II and future hadron colliders. The model is an extension of the SM that contains a color-triplet scalar, T, with quantum numbers $(\bar{3}, 1, 1/3)$. The corresponding Lagrangian is

$$\mathcal{L}_{\rm UV} = f_a\, T\, \bar{N}_a^c P_R b + f'\, T^*\, \bar{t}^c P_R b + M_T^2 |T|^2 \ . \tag{10}$$

In fact, this is the simplest model that yields the flavor

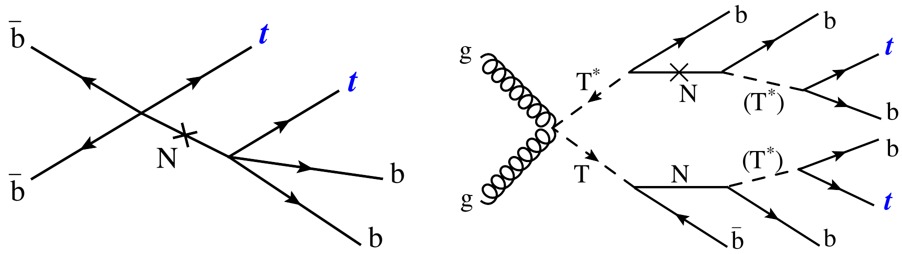

FIG. 3: Feynman diagrams for same-sign top quark events that can happen at hadron colliders, using the Majorana nature of RHNs and the baryonic interactions of Eq. (3). **Left**: $pp \to tN$ production via the contact operator Eq. (3), followed by the decay $N \to tbb$. **Right**: process in the UV complete model, pair production of color-triplet scalars T, T^*, and followed by $T \to N\bar{b}$, $T^* \to Nb$, and $N \to tbb$ (via virtual T in parentheses). Baryon number and top quark number are broken when both RHNs decay into top quarks using a Majorana mass insertion.

and color structures of the effective operators in Eq. (3), after integrating out the color-triplet scalar T, corresponding to a cutoff $\lambda_a/\Lambda^2 \equiv f_a f'/M_T^2$. The TeV scale cutoff as discussed above can be naturally obtained for $M_T \sim$ TeV and $f_a f' \sim \mathcal{O}(1)$. We note that in the above UV model, it is possible to have baryogenesis through the decays of the T particle [12]. We will not further explore such a possibility in this work.

The introduction of the scalar T could offer richer phenomenology at colliders. If light enough, T, T^* can be pair produced at hadron colliders. Each triplet will first decay into $N + b$, which is then followed by subsequent decay $N \to t\,b\,b$ via a virtual T. The above chain of processes are represented by the diagram in the right panel of Fig. 3. These together result in same-sign top quark final states with many b-jets. In Table II, we give the leading-order QCD cross section for the T, T^* pair production at the 13 TeV LHC and a 100 TeV proton-proton collider, calculated with MadGraph [32].

$\sigma(pp \to TT^*)$				
M_T	1.5 TeV	2 TeV	5 TeV	10 TeV
$\sqrt{s} = 13$ TeV	0.16 fb	0.01 fb	—	—
$\sqrt{s} = 100$ TeV	384 fb	92 fb	0.54 fb	4×10^{-3} fb

TABLE II: Pair production cross sections of T, T^* via strong interaction at the 13 and 100 TeV proton-proton colliders.

Moreover, an additional distinct signal could be displaced vertices from the decay of RHNs, if we take a somewhat larger cutoff scale $\Lambda/\sqrt{\lambda_a}$. In fact, we find for $M_1 = 200$ GeV and $\Lambda/\sqrt{\lambda_1} \gtrsim 7$ TeV, Eq. (4) implies a displaced decay length $c\tau_{N_1} \gtrsim 100$ μm, which would be detectable at the LHC [33]. This could result from Eq. (10) for $M_T \sim 1-2$ TeV and $f_a \sim f' \sim 0.2$. Events with same-sign tops and displaced vertices would be quite striking and hard to miss in collider experiments. Meanwhile, if the corresponding neutrino Yukawa coupling of N_1 is $y_N^1 \gtrsim 10^{-7}$, sufficient to explain the solar neutrino mass difference [27], the partial decay rate of $N_1 \to W\ell$ can be as large as order one. The leptonic decays can be used to identify N_1 as a RH neutrino (see, *e.g.*, [34]).

We note that the T particle introduced above has the same quantum numbers as the color-triplet partner of a Higgs doublet in the fundamental $\mathbf{5}_H$ representation of the SU(5) GUT [35]. This gives the motivation to consider a more unified framework for our scenario [36]. Our light color-triplet T cannot arise from the same $\mathbf{5}_H$ as the SM Higgs, whose Yukawa couplings are inconsistent with those in Eq. (10), within our framework. One could introduce a new $\bar{\mathbf{5}}'_H = (T, D)$ scalar, where D is a Higgs doublet with quantum numbers $(1, 2, -1/2)$ under the SM. The T couplings in Eq. (10) can then come from the SU(5) gauge invariant Yukawa interactions,

$$\begin{aligned} \mathcal{L}_{\mathrm{SU(5)}} &= f_a \mathbf{5}_3 \mathbf{1}_a \bar{\mathbf{5}}'_H + M_T^2 \bar{\mathbf{5}}'_H \mathbf{5}'_H \\ &+ f_0\, \overline{\mathbf{10}}_3 \mathbf{5}_3 \mathbf{5}'_H (\mathbf{24}_H + v_{\mathrm{GUT}})/\Lambda_{\mathrm{GUT}} \ , \end{aligned} \quad (11)$$

where $\mathbf{1}_a$ are RHNs and singlets under SU(5), the lower index 3 means only the third generation fermions are involved, and the bar over a representation means the complex conjugate of it. The last term contains a higher dimensional operator which after GUT symmetry breaking, $\langle \mathbf{24}_H \rangle = v_{\mathrm{GUT}}\, \mathrm{diag}(2/3, 2/3, 2/3, -1, -1)$, projects out the $\bar{t}^c P_R b T^*$ operator in Eq. (10) with $f' = 5 f_0 v_{GUT}/(3\Lambda_{\mathrm{GUT}})$. At the same time it forbids the dangerous operator $\bar{Q}LT$, thus evading the usual doublet-triplet splitting problem [37]. As a consequence of this setup, the quark Yukawa interaction $\bar{Q} b_R D$ of the SU(2) Higgs doublet D is also forbidden; it still possesses the neutrino Yukawa interaction from the f_a term. Without further splitting of the T and D components, the mass of D also lies at the TeV scale with a leptophilic nature.

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